Dark energy, the colored anti-de Sitter vacuum, and LHC phenomenology

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We study the possibility that the current accelerated expansion of the universe is driven by the vacuum energy density of a colored scalar field which is responsible for a phase transition in which the gauge $SU(3)_c$ symmetry breaks. We show that if we are stuck in a $SU(3)_c$-preserving false vacuum, then $SU(3)_c$ symmetry breaking can be accommodated without violating any experimental QCD bounds or bounds from cosmological observations. Moreover, unless there is an unnatural fine-tuning beyond the usual cosmological constant fine-tuning, the true vacuum state of the universe is anti-de Sitter. The model can likely be tested at the LHC. A possible (though not neccessary) consequence of the model is the existence of fractionally charged massive hadrons. The model can be embedded in supersymmetric theories where massive colored scalar fields appear naturally, thus reopening previously closed regions of SUSY parameter space.

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Observational data indicate that our universe is going through a phase of accelerated expansion. To date it remains a mystery what is the driving force behind the acceleration. Data favor an equation of state of the cosmic fluid $w \approx -1$, corresponding to a constant, or nearly constant, energy density. The null hypothesis is that we have reached the lowest energy state of the universe – the true vacuum energy density (or cosmological constant).

If it is indeed the true vacuum energy we are seeing, it may represent the worst prediction ever made by a theory. The value needed to explain the observed acceleration, $(10^{-3}eV)^4$, is 124 orders-of-magnitude smaller than the value naively predicted by particle theory, $(10^{19}GeV)^4$, and still 60-ish orders-of-magnitude after the assistance of supersymmetry.

It is equally plausible that the we are instead noticing the effects of some metastable false vacuum energy associated with some matter field. This would also have the equation of state $w = -1$. If this is the case, how might the true vacuum of the universe differ from the false vacuum in which we find ourselves? One startling realization is that just because the magnitude of the current vacuum energy is so close to zero, does not mean that the magnitude of the true vacuum energy is similarly fine-tuned. Indeed, it would in many ways be more natural for the true vacuum energy to be separated from ours by an amount characteristic of whatever phase transition is associated with the transition from true to false vacuum. The true vacuum state of the universe might therefore be strongly anti-de Sitter.

Non-zero vacuum energy density is usually associated with the breaking of some symmetry. What symmetry might be breaking? The most interesting possibility is that it is the unbroken gauge symmetry – $SU(3))_c \times U(1)_{\rm EM}$. The Universe in its history has apparently undergone a number of phase transitions. Currently, the physics is invariant under local $SU(3)_c \times U(1)_{\rm EM}$ transformations. However, there is no reason to believe that the chain of symmetry breaking stops here. We focus here on the possibility of $SU(3)_c$ breaking, though similar possibilities may exist for $U(1)_{EM}$ breaking.

The possible breaking of $SU(3)_c$ symmetry was intensively discussed some time ago for different reasons [1]. Depending on the model, the $SU(3)_c$ gauge symmetry can be broken completely or down to some subgroup of $SU(3)_c$. Regardless of the details, a generic feature of $SU(3)_c$ breaking is the existence of a new colored scalar Higgs multiplet, Φ. The non-zero vacuum expectation value of this field breaks the color symmetry and gives mass to gluons. To enforce the observed color symmetry locally, we will insist that the phase transition is first order, and that the observed universe is in a color-symmetric false vacuum.

Consider a Lagrangian density invariant under the $SU(3)_c$ transformations:

$$\mathcal{L} = -\frac{1}{2}{\rm Tr}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}{\rm Tr}(D_\mu\Phi)(D^\mu\Phi) - V(\Phi)\,. \quad (1)$$

For illustrative purposes, we take Φ in the adjoint representation of $SU(3)$. $D^\mu\Phi$ is the covariant derivative of Φ with a gauge coupling constant g. The potential

$$V(\Phi) = \frac{\mu^2}{4}({\rm Tr}\Phi^2) + \frac{\lambda_1}{16}({\rm Tr}\Phi^2)^2 + \frac{\lambda_2}{6}({\rm Tr}\Phi^3) + V_0\,. \quad (2)$$

A Hamiltonian bounded below is assured if $\lambda_1 > 0$. The term ${\rm Tr}\Phi^4$ is not included, but, since it has the same structure as $({\rm Tr}\Phi^2)^2$, it will not change our qualitative results. A similar model was studied in [7]. We include a constant V_0, which is an overall shift of the potential.

We choose to work in a diagonal representation where Φ is Hermitian and traceless. Let the three real fields of the diagonal representation be ψ_1, ψ_2 and $\psi_3 = -(\psi_1 + \psi_2)$. Minimizing V($\Phi$) we find

$$\psi_1 = \psi_2 = -\psi_3/2 \equiv \psi\,. \quad (3)$$

The potential as a function of ψ is:

$$V(\psi) = \frac{3}{2}\mu^2\psi^2 + \frac{9}{4}\lambda_1\psi^4 - \lambda_2\psi^3 + V_0\,. \quad (4)$$

Defining

$$\psi_0 \equiv \frac{2}{9}\frac{\lambda_2}{\lambda_1} \quad \text{and} \quad \epsilon_0 \equiv \lambda_1 - \frac{2\mu^2}{3\psi_0^2}\,, \quad (5)$$

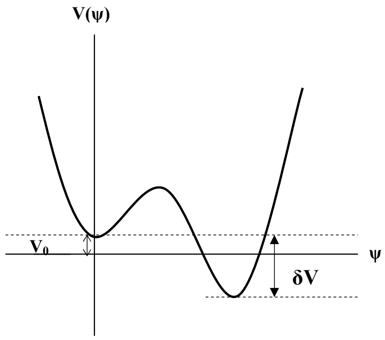

FIG. 1: Characteristic potential for a first order phase transition. There exist the false vacuum where the symmetry is unbroken, $\psi = 0$, and the true vacuum where the symmetry is broken, $\psi \neq 0$. r δV is the difference in energy densities between the vacua. V_0 plays the role of the cosmological constant.

allows us to rewrite $V(\psi)$ in a more convenient form:

$$V(\psi) = \frac{9}{4}\lambda_1\psi^2(\psi-\psi_0)^2 - \frac{9}{4}\epsilon_0\psi_0^2\psi^2 + V_0\,. \qquad (6)$$

The role of the parameter ϵ_0 is to introduce a controlled fine tuning in $V(\psi)$. If $\epsilon_0 = 0$, $V(\psi)$ has two degenerate minima at $\psi = 0, \psi_0$. However, for a small positive ϵ_0, the global minimum is at $\psi = \psi_0(1+\epsilon_0/\lambda_1)$. The difference between the energy densities of the two vacua is

$$\delta V = 9\epsilon_0\psi_0^4/4\,. \qquad (7)$$

We included V_0 in the potential so that we could specify the false vacuum energy density, $V(\psi=0) = V_0$. We require that V_0 is positive and extremely small, $V_0 \simeq (10^{-3}\text{eV})^4$, representing the vacuum energy density that is driving the accelerated expansion of the universe. This situation is depicted in Fig. 1. We do not explain the value of V_0 – it is an observational input, and its small value is the usual cosmological constant fine tuning, We will show, however, that that ϵ_0 does not have to be fine-tuned: $\epsilon_0 \lesssim 0.1$ will be sufficient to avoid all constraints. Thus, the shape of the potential is not fine-tuned, Since generically $\delta V \gg V_0$, the true vacuum energy density will be large and negative, representing an AdS space.

$V(\psi)$ is the zero-temperature potential. Since we are interested in the first-order phase transition between the false vacuum $\psi = 0$, and the true vacuum $\psi = \psi_0$, we need the finite-temperature effective potential. Following [7], we find that at one-loop

$$V_{\text{eff}}(\psi) = \frac{9}{4}\lambda_1\psi^2(\psi-\psi_0)^2 - \frac{9}{4}\epsilon(T)\psi_0^2\psi^2 + V_0\,, \qquad (8)$$

where

$$\epsilon(T) \equiv \epsilon_0 - \frac{2}{3}\frac{\mu^2(T)-\mu^2}{\psi_0^2} = \epsilon_0 - \left(\frac{5}{9}\lambda_1 + \frac{1}{2}g^2\right)\frac{T^2}{\psi_0^2}\,. \quad (9)$$

All temperature dependence is in $\mu^2(T) \equiv \mu^2 + (\frac{5}{6}\lambda_1 + \frac{3}{4}g^2)T^2$. The critical temperature at which $\epsilon(T) = 0$ is

$$T_c^2 = \frac{2\epsilon_0\psi_0^2}{10\lambda_1/9 + g^2}\,. \qquad (10)$$

For $T \gg T_c$, the potential has a single minimum at $\psi = 0$, color symmetry is thus unbroken. As T falls, a second minimum emerges at $\psi \neq 0$; this becomes the global minimum at $T < T_c$. Thus, the true vacuum state breaks $SU(3)_c$; however, our universe is stuck in a supercooled color-preserving false vacuum with $\psi = 0$. Thermal fluctuations at $T \ll T_c$ are not large enough to take the system over the barrier into the true vacuum.

The existence of a lower minimum indicates that any point in the universe will eventually tunnel into the true $SU(3)_c$-breaking vacuum. However, the tunneling rate depends on the energy difference between the vacua and the height of the potential barrier, and can be extremely slow as we discuss below. A universe in such a false vacuum will evolve as has ours – initially radiation dominated, then matter dominated, and eventually expanding at an accelerated rate as it becomes dominated by the false vacuum energy density, V_0.

In the false vacuum state, $\psi = 0$, excitations of the colored Higgs have mass $m_\psi \sim \mu$. It is natural to chose μ to be $\mathcal{O}(1\text{TeV})$ as this is the energy scale at which new physics beyond the Standard Model is expected. Without new physics at this scale, the Standard Model is severely fine-tuned. It is also very difficult to hide a colored scalar field much lighter that this. This choice automatically evades all of the experimental QCD constraints.

The energy density of the false vacuum must, as we noted, be chosen phenomenologically, $V_0 \equiv \rho_v \approx (10^{-3}\text{eV})^4$. However, an interesting numerology, $V_0 \approx \left(\text{TeV}^2/M_{\text{Pl}}\right)^4$, hints perhaps toward a gravitational origin of the overall shift of the potential if the fundamental scale of the physics is TeV.

We next examine the limits on the parameter ϵ_0. An important question is how long would the universe exist in the supercooled false vacuum state $\psi = 0$. The transition from this de Sitter-like false vacuum to the true anti-de Sitter-like vacuum, is realized by nucleation of bubbles of true vacuum inside the false vacuum. In the semi-classical approximation, the transition probability per unit space-time volume is

$$\Gamma = Be^{-S_E} \qquad (11)$$

where S_{E} is the Euclidean action of the $O(4)$-symmetric bounce solution describing the tunneling. B is a dimensionful constant $\mathcal{O}(\text{TeV}^4)$, the details of which we ignore. To calculate S_{E}, we follow the method developed in [2].

The one-dimensional, zeroth-order (in ϵ_0) Euclidean action per unit volume for the tunneling is

$$S_1 = \int_0^{\psi_0} d\psi' \sqrt{2V(\psi')} \approx \sqrt{\frac{\lambda_1}{8}}\psi_0^3\,. \qquad (12)$$

In the zero-temperature limit and thin wall approximation, the radius of the critical bubble is

$$R_0 = \frac{3S_1}{\delta V} = \frac{\sqrt{2\lambda_1}}{3}\frac{1}{\epsilon_0\psi_0}\,. \tag{13}$$

The Euclidean action for an $O(4)$ symmetric bubble is

$$S_E = -\frac{1}{2}\delta V\pi^2 R_0^4 + 2\pi^2 R_0^3 S_1 = \frac{\pi^2\lambda_1^2}{54}\frac{1}{\epsilon_0^3} \tag{14}$$

At zero temperature, the decay rate per unit volume per unit time is given by (11). In order for our observable universe, whose four-volume is of the order of $t^4_{\rm Hubble}$, to remain in the false vacuum, one needs roughly $\Gamma t^4_{\rm Hubble} \lesssim 1$. Taking $t_{\rm Hubble} \sim 10^{10}$years, we find that sufficient metastability of the color-neutral false vacuum is obtained for $S_{\rm E} \gtrsim 400$. With a generic value $\lambda_1 \sim 1$, we see that vacuum stability requires only $\epsilon \lesssim 0.1$, an extremely mild fine-tuning. On the other hand, $\epsilon \lesssim 0.1$ is enough to make the thin wall approximation valid.

To make sure that the story remains unchanged in the early universe, we calculate the temperature dependent decay rate in the high temperature limit. At finite temperature, instead of an $O(4)$ symmetric bounce, one should look for an $O(3)$ symmetric solution, periodic in time with period $1/T$ (see [8]). In the high temperature limit ($T >> 1/R_0$), the time integration in the calculation of the bounce action yields an overall factor of $1/T$. The decay exponent S_E now has the form

$$S_E = S_3(\psi, T)/T\,, \tag{15}$$

where S_3 is the three-dimensional action of the $O(3)$ symmetric bubble. The radius of the critical bubble is

$$R(T) = \frac{2S_1}{\delta V(T)} = \frac{\sqrt{8\lambda_1}}{9}\frac{1}{\epsilon(T)\psi_0}\,. \tag{16}$$

The three dimensional action of the bounce is given by

$$\begin{aligned} S_3(T) &= -\frac{4}{3}\pi R(T)^3\delta V(T) + 4\pi R(T)^2 S_1 \qquad (17)\\ &= \frac{\sqrt{128\lambda_1^3}}{243}\frac{\psi_0}{\epsilon^2(T)} \end{aligned}$$

The temperature dependent decay rate is now

$$\Gamma(T) \approx \exp\left[-\frac{\sqrt{128\lambda_1^3}}{243}\frac{\psi_0}{T\epsilon^2(T)}\right] \tag{18}$$

Thus, at high temperatures $T \gg T_c$, when the color preserving minimum at $\psi = 0$ is the lowest energy state, the transition rate is large and most of the universe ends up in that minimum. At $T = T_c$, where strictly speaking the high temperature approximation is not valid, at least formally $\epsilon(T) = 0$ and transitions between the vacua are suppressed. As shown earlier, for $T < T_c$, $\psi = 0$ is a false minimum, but the transition rate to the true color-breaking vacuum is suppressed by the bare value of ϵ_0. We saw that $\epsilon_0 \sim 0.1$ makes the transition time larger than the current Hubble time.

One might worry that, even at $T \ll T_c$, energetic processes in the history of our universe (e.g. cosmic ray collisions) would have stimulated the formation of a true vacuum bubble, and that such a bubble would have expanded rapidly to encompass most of the visible universe. Fortunately, it is not a simple thing to create a vacuum bubble in a high energy collision. This requires not just sufficient energy but a coherent superposition of a large number of high energy quanta over a volume large compared to the characteristic energy. Such processes happen freely at high temperature but essentially not at all at high energy. The height of the barrier between the false and true vacua is

$$V_{max} = 9\lambda_1\psi_0^4/64\,. \tag{19}$$

This is approximately TeV^4 for $\psi_0 \sim 1$TeV and $\lambda_1 \sim 1$. Such temperatures are unlikely to soon be achieved in colliders, and are probably not achieved over large enough volume even in the highest energy cosmic ray collisions.

The critical bubble radius (13) is 5TeV^{-1} for our canonical parameters ($\epsilon_0 \sim 0.1$, $\psi_0 \sim 1$TeV and $\lambda_1 \sim 1$). Given (19), this suggests that we need approximately

$$N_{quanta} \simeq \left(\frac{4\pi R_0^3}{3}V_{max}\right)V_{max}^{-1/4} \gtrsim 100\,(0.1/\epsilon)^3\,. \tag{20}$$

individual excitations coherently supposed. Note that N_{quanta} grow very fast – as ϵ^{-3}. To create a bubble in a high energy collision is thus extremely difficult. A bubble is a highly coherent state of Φ particles. Produce these quanta and they are more likely to fly apart than to assemble into a true vacuum bubble. The probability that you make a bubble should therefore be suppressed by a huge entropy factor. Cosmic rays, which usually produce collisions with $E_{cm} \lesssim 100$TeV, are thus extremely unlikely to trigger the destruction of the known universe.

The scenario proposed here should have distinct experimental signatures in near-future accelerators like the Large Hadron Collider (LHC) and in cosmic ray observatories. The principal generic feature of our scenario is the existence of colored scalar fields with mass around 1TeV. Experimental signatures depend crucially on the lifetimes of excitations of these fields, which are model dependent. A heavy colored scalar can generically decay into gluons and quarks. Its lifetime depends on couplings and on the specific representation of $SU(3)_c$. If the decay is fast, the main signature of its production will be gluon and quark jets in excess of the standard model prediction.

The lifetime of the colored scalar could be longer than the characteristic hadronization time $\sim$ (100MeV)$^{-1}$. This could happen if Φ is in a high-dimensional representation of $SU(3)_c$ or is protected by a symmetry, such as R-parity. In an unbroken phase, it is impossible to find an isolated colored particle; however, color singlets can be found free, thus Φ would combine into color singlet bound states. One possibility is bound state of two or more colored scalars (depending on the $SU(3)_c$ representation to which Φ belongs). It was argued in [5] that such states would resemble glueballs (bound states of gluons). Another accessible new state is a bound state of Φ with quarks. If the scalar particle does not carry electromagnetic charge, these states could be fractionally charged – a unique (though not generic) signature of our model. We would expect the lightest of these new states (e.g. a bound state of Φ with the lightest quarks) to have mass of the order of a few TeV.

Specific experimental signatures will depend on the exact model of $SU(3)_c$ breaking. For example, in the model discussed above, the symmetry breaking is realized by a scalar field that transforms under the eight-dimensional adjoint representation. Since 8×8 in the $SU(3)_c$ representation classification scheme contains an $SU(3)_c$ singlet state, one can find a color-free bound state of two scalar fields. Like glueballs, such states would dominantly decay into lighter mesons. Quarks (q) and antiquarks ($\bar{q}$) belong to 3 and $\bar{3}$ representations of $SU(3)_c$. The simplest combination that contains a color singlet is $8 \times 3 \times \bar{3}$, corresponding to $\Phi q\bar{q}$. This combination has integer charge. However, in models where $SU(3)_c$ is broken by a scalar multiplet belonging to a $3, 6, 10, 15$ or 21 dimensional representation, the bound states can have fractional charges. The 27 of $SU(3)$ again gives an electrically neutral bound state. Integer charged states decay to lighter mesons, whereas the lightest fractionally charged state is presumably stable. These experimental signatures will be explored in greater detail in future publications.

While discovery of a colored scalar field would be indicative, it would shed no light on the global shape of the potential. In [6], a method for reconstructing the potential was proposed. In a model containing kink solutions, one could in principle characterize the potential by studying the spectrum of kinks. Producing a kink corresponds to producing a small bubble of true vacuum. Since the potential barrier is only a TeV high, it might be possible to realize this in future experiments. Such bubbles would be subcritical, only TeV^{-1} in size, and would decay quickly. The decay pattern of subcritical bubbles would shed light on the global shape of the potential. Production of a large, supercritical bubble is suppressed. Thus, it might be possible to test the model without destroying the known universe.

Finally we mention the possibility that an $SU(3)_c$-breaking true vacuum can be embedded into supersymmetric models. Supersymmetric models naturally contain colored massive scalar fields as super-partners of the ordinary standard model quarks. It has been thought that such fields are dangerous since they can lead to breaking of the $SU(3)_c$ symmetry, *i.e.* the true minimum in the landscape of supersymmetric minima may be one with broken $SU(3)_c$. A large portion of the supersymmetric parameter space is therefore excluded [4]. Based on the discussion presented here, we propose that one does not need to do this. In particular, false minima where $SU(3)_c$ is not broken and which lie above the true vacuum can drive the accelerated expansion of our universe (for related work see [9]). Note that in this case Φ is likely a squark field and so carries fractional electric charge. Φ cannot however be the dark matter (nor can the dark matter be a bound state involving Φ, as that would contradict well-known limits on strongly interacting dark matter [10]).

In conclusion, we have proposed that our universe is stuck in a false vacuum where $SU(3)_c$ symmetry remains unbroken, while the true vacuum breaks $SU(3)_c$. The vacuum energy density of the false vacuum drives the observed accelerated expansion. A slight tuning of the potential, $\epsilon_0 \sim 0.1$, makes the transition time to the true vacuum larger than the Hubble time. Massive colored scalar fields provide unique experimental signatures in near-future accelerator experiments. Discovery of fractionally charged hadrons would be particularly strong evidence for our model. Using recently developed techniques, one could in principle determine the shape of the potential by solving the inverse scattering problem. We suggest that a large portion of the SUSY parameter space has been incorrectly excluded. We also note that with small modifications, a similar mechanism can be applied to $U(1)_{\text{EM}}$, the gauge group of electromagnetism or to the conventional Higgs of the Standard Model.

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